

Extremely efficient Zevatron in rotating AGN magnetospheres.

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ABSTRACT

A novel model of particle acceleration in the magnetospheres of rotating active galactic nuclei (AGN) is constructed. The particle energies may be boosted up to 10^{21} eV in a two step mechanism: In the first stage, the Langmuir waves are centrifugally excited and amplified by means of a parametric process that efficiently pumps rotational energy to excite electrostatic fields. In the second stage, the electrostatic energy is transferred to particle kinetic energy via Landau damping made possible by rapid "Langmuir collapse". The time-scale for parametric pumping of Langmuir waves turns out to be small compared to the kinematic time-scale, indicating high efficiency of the first process. The second process of "Langmuir collapse" - the creation of caverns or low density regions - also happens rapidly for the characteristic parameters of the AGN magnetosphere. The Langmuir collapse creates appropriate conditions for transferring electric energy to boost up already high particle energies to much higher values. It is further shown that various energy loss mechanisms are relatively weak, and do not impose any significant constraints on maximum achievable energies.

Key words: acceleration of particles - cosmic rays - galaxies: active

1 INTRODUCTION

New observations confirming a strong correlation of ultra-high energy cosmic ray protons with the active galactic nuclei (AGN) (Kim & Kim 2013) has stimulated the search for a possible mechanism that could accelerate protons up to tens of EeV and higher energies.

In this paper we will explore energy transfer via the relativistic centrifugal force in the rotating magnetospheres of compact objects as a mechanism for boosting proton energies to such enormous levels. In general, this is not a new idea, Bogovalov & Tsinganos (1999) and Bogovalov (2001) have considered the role of magnetocentrifugal effects in astrophysical outflows and nonaxisymmetric magnetospheres. Based on a rather different approach the mechanism of centrifugal acceleration was applied to explain the observed high and very high energy emission pattern (Osmanov et al. 2007; Osmanov 2010) in some astrophysical settings.

Relativistic centrifugal force turns out to be a crucial physical element governing particle acceleration, particularly, in astrophysical objects endowed with rotating magnetospheres, where magnetic energy density is larger than that of the plasma. The plasma particles are, thus, forced

to follow the field lines, which in turn are corotating. Consequently, the particles, sliding along these magnetic field lines, experience centrifugal force and gain energy.

The time-varying centrifugal force can also parametrically excite electrostatic modes providing an additional channel for transferring energy from the rotator to the electromagnetic waves and, then, possibly to the particles; this channel may be particularly relevant to the enormously energetic AGN magnetosphere (Machabeli et al. 2005; Osmanov 2008).

The energy pumping mechanism has already been applied to both pulsars (Machabeli et al. 2005) and the AGN (Osmanov 2008). In the AGN magnetosphere, the centrifugally excited electrostatic waves are found to be unstable; the growth rates, exceeding the corresponding rate of the accretion disc evolution, indicate the possibility of an efficient instability (Osmanov 2008).

Till now, our most extensive exploration of the parametrically driven Langmuir waves and their role in accelerating particles has been in the context of the Crab nebula (Machabeli et al. 2005; Mahajan et al. 2013). We showed that the process of generation of the centrifugally driven Langmuir waves is rather efficient, and for the Crab pulsar magnetosphere, energy gain of electrostatic modes might be in a reasonable coincidence with its observed slow down

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rate. These waves, sustained by the main electron-positron (magnetospheric) plasma, were shown to Landau damp effectively on the primary electron beam (much less density than the main plasma), imparting large per particle energy to the already fast particles. We have argued that the Crab pulsar, with its extremely high rate of rotation, may guarantee acceleration of electrons to hundreds of TeV and even higher.

In this paper, we go back to reexamine the problem begun in (Osmanov 2008); we will investigate the efficiency of centrifugally driven electrostatic waves in the AGN magnetosphere, and determine if these waves can be a conduit for transferring energy towards particle acceleration. In contrast to the pulsars, the AGN atmosphere contains a quasi-neutral plasmas composed of protons and electrons. Consequently the mass dependent centrifugal force acting will produce an initial charge separation on electrons and protons, leading, perhaps, by means of the parametric instability, to extremely efficient energy pumping from the rotating magnetosphere. In addition to demonstrating the existence of unstable Langmuir waves, we also explore a mechanism that might be responsible for particle acceleration.

It will be helpful to remember that the Langmuir turbulence inevitably leads to (nonlinear) instability, creating caverns-low density regions (Zakharov 1972). According to the investigation, the high frequency contribution of pressure inside the cavern pulls the particles from this area, provoking an explosive collapse of the cavern, efficiently transferring energy from the electrostatic waves to the particles pushed from the collapsing regions. Dynamics of collapse was studied in detail in (Galeev et al. 1977) and numerically investigated by Degtiarev et al. (1976). Galeev et al. (1977) considered a theory of three dimensional problem and investigated the corresponding spectra of Langmuir turbulence, studying the role of absorption mechanisms for short wavelength plasmons. Degtiarev et al. (1976) analysed the dynamics of modulational instability for long-wavelength electrostatic oscillations and numerically studied the role of Langmuir damping in the termination of collapse.

After spelling out our theoretical model in Sec.2, we work out in Sec.3 the details of the particle acceleration mechanism for typical parameters of AGN magnetospheres. In Sec. 4, we discuss and summarize our results.

2 THEORETICAL MODEL

Transferring the approach of Mahajan et al. (2013) to the AGN atmosphere, we divide the problem into two subtasks: (a) generation of electrostatic waves, and (b) Particle acceleration which, in this case, happens through the Langmuir collapse.

2.1 Centrifugal excitation of Langmuir waves

For the mode calculations of this paper, the magnetic field lines are almost straight and co-rotate with the supermassive black hole. For a typical AGN with mass $M \sim 10^8 \times M_\odot$, the angular velocity of rotation is given by

$$\Omega \approx \frac{ac^3}{GM} \approx 10^{-3} \frac{a}{M_8} \text{rad/s}^2, \quad (1)$$

where $M_8 \equiv M/(10^8 M_\odot)$ and $0 < a \leq 1$ is a dimensionless parameter characterizing the rate of rotation of the black hole. The magnetic field may be estimated by assuming that the magnetic energy density is of the order of radiation energy density of the AGN (equipartition) (Osmanov et al. 2007),

$$B \approx \sqrt{\frac{2L}{r^2 c}} \approx 27.5 \times \left(\frac{L}{10^{42} \text{erg/s}} \right)^{1/2} \times \frac{R_{lc}}{r} G, \quad (2)$$

where L is the bolometric luminosity of AGN, $R_{lc} = c/\Omega$ is the light cylinder radius and r is the distance from the black hole. With these fields, the protons with high Lorentz factors, $\gamma \sim 10^{4-5}$, will have gyroradii much smaller than the kinematic length-scale (R_{lc}) of particles in the magnetosphere. Such particles will, clearly, follow and co-rotate with the magnetic field lines. The protons will, therefore, inevitably accelerate centrifugally, but as it is explained in detail by Osmanov et al. (2007), this process is strongly limited by two major factors.

In due course of motion protons gain energy, but this process lasts until the energy gain is balanced by energy losses due to the inverse Compton scattering. The maximum attainable proton Lorentz factor is, then, estimated as (Osmanov et al. 2007)

$$\gamma_{IC} \approx \left(\frac{6\pi m_p c^4}{\sigma_T L \Omega} \right)^2 \approx 1.5 \times 10^{11} \left(\frac{10^{42} \text{erg/s}}{L} \right)^2, \quad (3)$$

where $m_p \approx 1.67 \times 10^{-24} \text{g}$ is proton's mass and $\sigma_T \approx 6.65 \times 10^{-25} \text{cm}^2$ is the Thomson cross-section. We have taken into account that the Black hole rotates with 10% of its maximum rate ($a = 0.1$).

This estimate, however, turns out to be a bit too optimistic. A more restrictive limit is set by the proton dynamics as it moves under the combined influence of the Coriolis and Lorentz forces. This limiting mechanism, called, for want of a better name, "breakdown of the bead on the wire (BBW) approximation (Rieger & Mannheim 2000)", yields a more moderate gamma (see (Osmanov et al. 2007) for details)

$$\gamma_{BBW} \approx \frac{1}{c} \left(\frac{e^2 L}{2m_p} \right)^{1/3} \approx 1.4 \times 10^5 \left(\frac{L}{10^{42} \text{erg/s}} \right)^{1/3}, \quad (4)$$

From Eqs. (3,4), it is clear that $\gamma_{BBW} \ll \gamma_{IC}$. BBW, thus, sets the theoretical upper limit on the Lorentz factor $\sim 1.4 \times 10^5$. Via centrifugal acceleration, the energy is pumped from the rotating magnetosphere. Although the centrifugally accelerated protons are strongly limited in energy, the centrifugal mechanism of energy pumping is very much in action.

To study the generation of extremely unstable Langmuir waves, we follow the method developed in (Machabeli et al. 2005; Osmanov 2008). Throughout the paper, we assume that the magnetic field lines are almost straight and lie in the equatorial plane. Then, in the 1 + 1 formalism (Thorne et al. 1986), the Euler equation in the corotating frame can be written as

$$\frac{d\mathbf{p}_\beta}{d\tau} = \gamma_\beta \mathbf{g} + \frac{e_\beta}{m_\beta} \left(\mathbf{E} + \frac{1}{c} \mathbf{v}_\beta \times \mathbf{B} \right), \quad (5)$$

β denotes the species index (electrons and protons), \mathbf{p}_β and \mathbf{V}_β are, respectively, the dimensionless momentum

($\mathbf{p}_\beta \rightarrow \mathbf{p}_\beta/m_\beta$) and velocity, $\mathbf{g} \equiv -\nabla\xi/\xi$ is the gravitational potential (a crucial component of the model), $e_\beta(m_\beta)$ is the charge(mass) of the corresponding particle, $d\tau \equiv \xi dt$, $\xi \equiv (1 - \Omega^2 r^2/c^2)^{1/2}$ and $\gamma_\beta \equiv (1 - \mathbf{V}_\beta^2/c^2)^{-1/2}$ is the Lorentz factor.

Transition to the laboratory frame (LF) is accomplished through the identity $d/d\tau \equiv \partial/(\xi\partial t) + (\mathbf{v}_\beta \cdot \nabla)$, and the relation $\gamma = \xi\gamma'$ connecting the corotating and LFs of reference (the prime denotes the quantity in the LF). The LF equation of motion

$$\begin{aligned} & \frac{\partial \mathbf{p}_\beta}{\partial t} + (\mathbf{v}_\beta \cdot \nabla) \mathbf{p}_\beta = \\ & = -c^2 \gamma_\beta \xi \nabla \xi + \frac{e_\beta}{m_\beta} \left(\mathbf{E} + \frac{1}{c} \mathbf{v}_\beta \times \mathbf{B} \right), \end{aligned} \quad (6)$$

coupled with the continuity

$$\frac{\partial n_\beta}{\partial t} + \nabla \cdot (n_\beta \mathbf{v}_\beta) = 0, \quad (7)$$

and the Poisson equation,

$$\nabla \cdot \mathbf{E} = 4\pi \sum_\beta e_\beta n_\beta. \quad (8)$$

completes the system.

In the equilibrium, the plasma is assumed to obey the frozen-in condition $\mathbf{E}_0 + \frac{1}{c} \mathbf{v}_{0\beta} \times \mathbf{B}_0 = \mathbf{0}$. The corresponding trajectory of the centrifugally accelerated particles, calculated by the standard single particle approach, turns out to be: $r_\beta(t) \approx \frac{V_{0\beta}}{\Omega} \sin(\Omega t + \phi_\beta)$ (see Appendix, Eq. (A7)) and $v_{0\beta}(t) \approx V_{0\beta} \cos(\Omega t + \phi_\beta)$, where we have taken into account the specific initial phases of particles, ϕ_β . It is worth noting that in spite of the aforementioned behaviour of the radial coordinate, it does not mean that the particle will oscillate. As we will see later, the instability becomes so efficient that the corresponding time-scale of energy conversion will be less than the rotation period.

Since different species experience different forces, the Euler equation will lead to spacial separation of charges, which in turn, generates an electrostatic field via the Poisson equation; under certain conditions, these fields may grow in time. For studying the development of Langmuir waves, we expand all physical quantities around this equilibrium state (keeping up to linear terms in perturbations):

$$\Psi \approx \Psi^0 + \Psi^1, \quad (9)$$

where $\Psi = \{n, \mathbf{v}, \mathbf{p}, \mathbf{E}, \mathbf{B}\}$. Fourier decomposing the perturbations,

$$\Psi^1(t, \mathbf{r}) \propto \Psi^1(t) \exp[i\mathbf{k}\mathbf{r}], \quad (10)$$

one obtains the linearized system of equations (6-7)

$$\frac{\partial p_\beta}{\partial t} + ikv_{\beta 0} p_\beta = v_{\beta 0} \Omega^2 r_\beta p_\beta + \frac{e_\beta}{m_\beta} E, \quad (11)$$

$$\frac{\partial n_\beta}{\partial t} + ikv_{\beta 0} n_\beta + ikn_{\beta 0} v_\beta = 0 \quad (12)$$

$$ikE = 4\pi \sum_\beta n_{\beta 0} e_\beta, \quad (13)$$

describing the evolution of the electrostatic field.

In terms of an effective density, defined by

$$n_\beta = N_\beta e^{-\frac{iV_{\beta 0} k}{\Omega} \sin(\Omega t + \phi_\beta)}, \quad (14)$$

one can derive, after straightforward algebra, the following set of non-autonomous "Mode" equations (Mahajan et al. 2013)

$$\frac{d^2 N_p}{dt^2} + \omega_p^2 N_p = -\omega_p^2 N_e e^{i\chi}, \quad (15)$$

$$\frac{d^2 N_e}{dt^2} + \omega_e^2 N_e = -\omega_e^2 N_p e^{-i\chi}, \quad (16)$$

where $\chi = b \cos(\Omega t + \phi_+)$, $b = \frac{2ck}{\Omega} \sin \phi_-$, $2\phi_\pm = \phi_p \pm \phi_e$ and $\omega_{e,p} \equiv \sqrt{4\pi e^2 n_{e,p}/m\gamma_{e,p}^3}$ and $\gamma_{e,p}$ are the relativistic plasma frequencies and the Lorentz factors for the stream components. Note that for simplicity, we have assumed the particle population to consist of only two streams with initial phases ϕ_\pm .

After Fourier transforming Eqs. (15,16) in time, we derive the "dispersion relation" of the corresponding modes

$$\omega^2 - \omega_e^2 - \omega_p^2 J_0^2(b) = \omega_p^2 \sum_\mu J_\mu^2(b) \frac{\omega^2}{(\omega - \mu\Omega)^2}, \quad (17)$$

where $J_\mu(x)$ is the Bessel function. Referring the reader to (Mahajan et al. 2013) for details, we just state here that this system may undergo an instability when the real part of the frequency satisfies the resonance condition, $\omega_r = \mu_{res}\Omega$. After expressing the frequency, $\omega = \omega_r + \Delta$ one can straightforwardly reduce the above equation

$$\Delta^3 = \frac{\omega_r \omega_p^2 J_{\mu_{res}}(b)^2}{2}, \quad (18)$$

that has a pair of complex solutions implying a growth rate (imaginary part of Δ)

$$\Gamma = \frac{\sqrt{3}}{2} \left(\frac{\omega_e \omega_p^2}{2} \right)^{\frac{1}{3}} J_{\mu_{res}}(b)^{\frac{2}{3}}, \quad (19)$$

where $\omega_r = \omega_e$. The time dependent centrifugal force that parametrically drives the electrostatic waves is different for the two species- so are their Lorentz factors.

Due to the presence of the Bessel function, which for large values of the index ($\mu_{res} = \omega_r/\Omega \gg 1$ tends to be nonzero only when the argument b is of the order of the index, the growth rate will peak when $b = 2ck \sin \phi_-/\Omega = \mu_{res} = \omega_r/\Omega$. For these most unstable modes, the phase velocity $v_{ph} \equiv \omega/k = 2c \sin(\phi_-)$ will exceed the speed of light for certain values of ϕ_- . Since there are no particles with such velocities, such waves will not Landau damp on the particles. However, if there is a process that can enhance the wave vector, Landau damping might be restored. It is in this context that we discuss, in the next subsection, the possibility of Langmuir collapse. The length scale of a "cavern" might significantly decrease, leading to a decrease of phase velocity, so that there are enough resonant particles for drawing energy from the electrostatic field.

2.2 Acceleration mechanism

It is worth noting that due to a small initial amplitude of the Langmuir wave, the high frequency pressure increases, pushing out the particles from the perturbed zone. The resulting polarization creates an additional electrostatic field which causes a further decrease in density in this region, creating what are termed as "caverns". Simultaneously, the plasmons (quasiparticles) accelerate towards these cavities, enter

them and enlarge their depth even more. The boosted high-frequency pressure induces an auto modulation instability of the spatial distribution of plasmons.

For the purpose of this paper, we will consider a kinematically relativistic fluid with non-relativistic temperatures. In the rest frame, it has been shown that the fluid obeys the nonlinear set of hydrodynamic equations (Zakharov 1972),

$$\frac{\partial}{\partial t} \delta n + n_0 \nabla \mathbf{v}_e = 0, \quad (20)$$

$$\frac{\partial}{\partial t} \mathbf{v}_e + \frac{e}{m} \nabla \varphi_e + \frac{3}{2} \frac{T_e}{2mn_0} \nabla \delta n = 0, \quad (21)$$

reduces to

$$\left[\frac{\partial^2}{\partial t^2} - 3\lambda_D^2 \omega_p^2 \frac{\partial^2}{\partial x^2} - \omega_p^2 \right] E = \frac{\delta n}{n_0} \omega_p^2 E, \quad (22)$$

$$\left[\frac{\partial^2}{\partial t^2} - \lambda_D^2 \omega_p^2 \frac{\partial^2}{\partial x^2} \right] \delta n = \frac{1}{16\pi m_p} \frac{\partial^2 E^2}{\partial x^2}, \quad (23)$$

where δn is the electron density perturbation, \mathbf{v}_e is the corresponding velocity perturbation, n_0 is the unperturbed ion number density, T_e is the electron temperature, φ_e is the high frequency part of the electrostatic potential and $\lambda_D \equiv \sqrt{T_e/(4\pi n_0 e^2)}$ is the Debye length scale. We assume that the plasma is quasi neutral.

It has been shown that this system, generalized for higher dimensions, describing the Langmuir collapse, exhibits explosive behaviour, i.e. the initial growth rate of the instability (Artsimovich & Sagdeev 1979)

$$\Gamma_{LC} = Im(\omega) \approx \frac{1}{\gamma_p} \left[\frac{\langle E^2 \rangle e^2 m_p}{4k_B T} \right]^{1/2}, \quad (24)$$

since it scales with the field intensity, will naturally increase with time. In deriving (24), we have taken into account the Lorentz boosting. According to the standard approach developed by Zakharov (1972), kinetic and potential energies of the plasmons caught by the caverns are of the same order of magnitude

$$k^2 \lambda_D^2 \sim \frac{|\delta n|}{n_0}. \quad (25)$$

Thus the characteristic length-scale of the cavern L_c , defined by the wavelength of the plasmon ($1/k$), scales inversely with density perturbation, i.e. $L_c \propto 1/\sqrt{\delta n}$. Since the high frequency pressure, $P_{hf} = -|E|^2 \delta n / (24k^2 \lambda_D^2 n_0)$ (Artsimovich & Sagdeev 1979) with $|\delta n|$, the greater the pressure the lesser the cavern width. This is precisely the recipe for an explosive instability leading to a smaller and smaller L_c . During this process, energy density of oscillations drastically increases accelerating the collapse. What happens later depends strongly on dimensionality of the process. If the collapse continues so that the $L_c = 1/k$ when the wavelength of plasmons is of the order of the Debye scale, the resonance Landau damping becomes important, resulting in particle acceleration. More discussion follows in the next section.

3 DISCUSSION

In this section we consider the implications of the theoretical framework developed in the last section (instabil-

ity of centrifugally induced electrostatic waves, and the expected Langmuir collapse) for particle acceleration in a typical AGN setting.

It has been assumed that the magnetic field are so strong that inside the light cylinder zone the AGN magnetospheric plasma co-rotates rigidly. Therefore, in this area the plasma number density might be well approximated by the Goldreich-Julian density

$$n_0 = \frac{\Omega B}{2\pi e c}. \quad (26)$$

If we assume that the maximum electron Lorentz factor is controlled by the same mechanism as protons ($\gamma_p \sim 1.4 \times 10^5$), then their lighter mass will allow $\gamma_e \approx \gamma_p (m_p/m_e)^{1/3} \sim 1.6 \times 10^6$ (see equation (4)) to be an order of magnitude larger. Considering two representative streams with Lorentz factors $\gamma_1 \sim 1.6 \times 10^6$ and $\gamma_2 \sim 10^3$ and temperature $T \sim 10^4$ K, one estimates that the instability time-scale, $1/\Gamma$, is less than the kinematic time-scale, $\sim 2\pi/\Omega$, implying a rather efficient process of pumping rotation energy into Langmuir waves. These waves will, then, accelerate particles by Landau damping aided by a possible Langmuir collapse.

Since the perturbation of density modulation usually is small, $\delta n \ll n_0$, the change of frequency of plasmons is negligible, $\delta \omega \ll \omega$. One may assume, then, that the corresponding energy is constant (Artsimovich & Sagdeev 1979)

$$\int d^q \mathbf{r} |E|^2 = const. \quad (27)$$

implying the scaling

$$|E|^2 \propto \frac{1}{l^q}, \quad (28)$$

where q equals $\{1, 2, 3\}$ depending on the dimension of the space. Inside the magnetosphere the magnetic field is strong enough to force particles to move along the field lines. The relevant geometry, thus, is 1 D and the high frequency pressure, $P \propto |E|^2$ scales as $1/l$ (Artsimovich & Sagdeev 1979). On the other hand, the thermal pressure, $P_{th} = k_B T \delta n$, behaving as $1/l^2$ (see equation (25)) increases faster than the high frequency pressure. This means that inside the magnetosphere the centrifugally induced Langmuir waves do not collapse and only propagate towards the outer region of the magnetosphere.

Outside the magnetosphere, the plasma kinematics are no longer governed by rotation. In this region, the density is controlled by the accretion processes. To estimate the density let us assume a spherically symmetric accretion. The expected accretion rate of particles per unit time and unit area is of the order of nv , where n is the accretion particle number density and $v = \sqrt{GM_{BH}/R_{lc}}$. The estimated density, then, is

$$n = \frac{L}{4\eta\pi m_p c^2 v R_{lc}^2} \approx 6.3 \times 10^5 \times \left(\frac{L}{10^{42} \text{ erg/s}} \right) \text{ cm}^{-3}. \quad (29)$$

We have assumed that only 10% of the rest energy of accretion matter ($\eta = 0.1$) transforms to emission. For such dense matter the corresponding plasma frequency, $\omega_p = \sqrt{4\pi e^2 n/m_p}$ exceeds the cyclotron frequency for protons $\omega_B = eB_{lc}/(m_p c)$, which means that the particles in the region outside the magnetosphere are no longer bound by

the magnetic field and consequently $d = 3$. This in turn means that the high frequency pressure behaving as $1/l^3$, increases much faster than the thermal pressure, behaving as $1/l^2$ and correspondingly the collapse of the Langmuir waves becomes inevitable.

During the collapse the density perturbation satisfies the approximate equation (Eqs. (22,25))

$$\frac{\partial^2 \delta n}{\partial t^2} \approx \frac{\delta n |E|^2}{16\pi n m_p \lambda_D^2}, \quad (30)$$

where the thermal pressure has been neglected.

After complementing this equation with already discussed relations, $|E|^2 \sim 1/l^3$ and $\delta n \sim 1/l^2$, one obtains (Zakharov 1972)

$$|E|^2 \approx |E_0|^2 \left(\frac{t_0}{t_0 - t} \right)^2 \quad (31)$$

$$l \approx l_0 \left(\frac{t_0}{t_0 - t} \right)^{-2/3}, \quad (32)$$

where t_0 is the time when the cavern collapses completely.

It is worth noting that the initial Langmuir waves have been efficiently amplified in the very vicinity of the light cylinder surface. The corresponding lengthscale can be obtained by a simple approximate expression $\Delta r \approx \gamma/(d\gamma/dr)$, leading to

$$\Delta r \approx \frac{\gamma_0}{2\gamma} R_{lc}, \quad (33)$$

where we have taken into account the radial behaviour of Lorentz factors of centrifugally accelerated particles (Rieger & Mannheim 2000)

$$\gamma(r) = \frac{\gamma_0}{1 - \frac{r^2}{R_{lc}^2}}. \quad (34)$$

As we have already discussed the electrostatic field appears and amplifies by means of the separation of charges in the mentioned zone. Therefore, the electrostatic field is approximated by the Poisson equation

$$E_0 \approx 4\pi n e \Delta r, \quad (35)$$

which due to the Langmuir collapse will be boosted by the factor of $(\Delta r/l)^{3/2}$ (see Eqs.(31,32)) where $l \approx 2\pi\lambda_D$ is the dissipation lengthscale (Artsimovich & Sagdeev 1979). It is clear that at the final stage, energy of the amplified electrostatic field will transfer to the particles inside the cavern resulting in protons with extremely high energies:

$$\epsilon_p \approx \frac{E^2}{8\pi n} = \frac{n e^2}{4\pi^2 \lambda_D^3} \Delta r^5. \quad (36)$$

For a proton beam with $\gamma_p \sim 10^2$, one can easily calculate that even the initial instability time-scale measured in the lab frame $\sim 1/\Gamma_{LC} \sim 0.2$ s is several orders of magnitude less than the kinematic time-scale. This difference will significantly increase, because the collapse has an explosive character and the electric field amplifies faster than the linear exponential increase. The Langmuir collapse is strong enough to guarantee efficient acceleration of particles to ultra high energies. In particular, from equation (36), one obtains

$$\epsilon_p \text{ (eV)} \approx 6.4 \times 10^{17} \times \left(\frac{f}{10^{-3}} \right)^3 \times \left(\frac{10^2}{\gamma_2} \right)^5 \times M_8^{-5/2} \times L_{42}^{5/2}, \quad (37)$$

where $f = \delta n/n_0$ is the initial dimensionless density perturbation, $a = 0.1$, $\eta = 0.1$, $L_{42} \equiv L/10^{42} \text{ erg/s}$ and we have assumed $T \sim 10^4 \text{ K}$. As it is evident from this expression for a convenient set of parameters one can achieve enormous energies of the order of 10^{21} eV . To achieve such energies, the required electrostatic fields must exceed the background magnetic field by many orders of magnitude (see Eqs. (2,35,36)). This is possible because the origin of the electrostatic field is different from that of the background magnetic field; the magnetospheric rotation energy is almost a limitless and continuous source, and can readily feed electric fields of such enormous magnitude. For the parameters corresponding to the maximum attainable energy 10^{21} eV , the length-scale of the cavern is of the order of 10^{12} cm (see equation 33). Again such a large scale structure of the electrostatic field can be maintained only because the available energy budget that is transferred to the Langmuir modes, is huge. It should also be stressed that the equilibrium frozen in condition (a condition relating the large scale equilibrium fields E_0 and B_0) is not affected by the much shorter scale electric fields associated with the Langmuir wave.

During the acceleration phase, the particles might lose energy due to several mechanisms that, potentially, might reduce the overall efficiency of acceleration. Via the highly efficient synchrotron mechanism, for example, the particles may rapidly lose their perpendicular momentum, transit to the ground Landau level and slide along the magnetic field lines. This mechanism, thus, does not influence particle acceleration.

The Inverse Compton scattering is also found to be little significance to the acceleration process. For such high energies, the relevant regime is Klein-Nishina, and the corresponding cooling time-scale goes as $t_{IC} = \epsilon_p/P_{KN} \propto \epsilon_p$ (Osmanov & Rieger 2009), where P_{KN} is the power emitted per unit time; P_{KN} is not sensitive to ϵ_p (Blumenthal & Gould 1970). Since t_{IC} is a continuously increasing function of ϵ_p , the inverse Compton process does not impose any constraints on achievable particle energies.

Another mechanism is the curvature radiation, characterized by the cooling time-scale $t_{cur} = \epsilon_p/P_{cur}$, where $P_{cur} = 2e^2 \epsilon_p^4 / (3m_p^4 c^3 \rho)$ is the energy loss rate and ρ is the curvature radius of the trajectory of particles. Acceleration is efficient until the acceleration time-scale, which is the collapse time-scale, is less than the cooling time-scale. Maximum energy is achieved when the following condition $t_{col} \approx t_{cur}$ is satisfied. If one takes into account the gyro radius, R_p , of relativistic protons and assumes $\rho \sim R_p$, one can obtain $\epsilon_{max} \approx 4 \times 10^{24} \xi \text{ (eV)}$, where $\xi \gg 1$ and characterizes the fact that the growth rate is much higher than the initial increment (see equation (24)). Therefore, the curvature emission is also negligible and cannot impose notable constraints on the maximum energies of particles.

Although we see that the present mechanism, in principle, might create primary cosmic ray (proton) energies in the ZeV range, the actual energies may be limited due to the interaction of these particles with the isotropic microwave cosmic radiation. This interaction could significantly reduce the proton energy to the so called GZK limit $4 \times 10^{19} \text{ eV}$ (Zatsepin & Kuz'min 1966; Greisen 1966). There is, however, ample observational evidence of cosmic rays with energies above the GZK limit (Abraham et al. 2008). This can happen, for instance, if there is not enough time for the

background radiation to slow down the more energetic particles - that is- the source of highly energetic particles lies within what may be called the the GZK radius, $\sim 100\text{Mpc}$ (Kim & Kim 2013)- the typical distance needed for significant energy loss on the cosmic microwave photons.

In the context of this paper it is worth noting that, the strong correlation of AGN with cosmic rays is actively discussed by Kim & Kim (2013), who consider the possibility that a subclass of AGN, might be responsible for ultra-high energy cosmic rays.

4 SUMMARY

(i) We have developed a new mechanism of particle acceleration operating in AGN magnetospheres. The mechanism, capable of creating extremely high energy cosmic rays, consists of two major stages: (I) centrifugal excitation of Langmuir waves and (II) the collapse of these waves by means of the modulational instability leading to particle acceleration.

(ii) For studying the excitation of electrostatic waves parametrically induced by relativistic centrifugal force we considered the linearized system of equations governing this process. It was found that the growth time of the instability is small compared to the kinematic time-scale, indicating the high efficiency of energy pumping from rotation to Langmuir waves in the light cylinder zone, where relativistic centrifugal effects become important.

(iii) As a next step we examined the Langmuir collapse of caverns (low density regions) by means of the high frequency pressure, exceeding the thermal pressure. This in turn, results in acceleration of protons up to ZeV energies. It has also been found that during acceleration major mechanisms governing energy losses do not impose any significant constraints on the maximum attainable proton energies.

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APPENDIX A:

In this section we consider a single particle, sliding along a corotating straight wire. It is clear that in the LF of reference the particle experiences a reaction force from the wire. Determining the relativistic momentum of particles

$$P_r = \gamma m \frac{dr}{dt}, \quad (\text{A1})$$

$$P_\varphi = \gamma m r \Omega, \quad (\text{A2})$$

where $\gamma = (1 - \Omega^2 r^2/c^2 - v^2/c^2)^{-1/2}$ is the Lorentz factor and $v \equiv dr/dt$, the equations of motion in the LF of reference are given by

$$\frac{dP_r}{dt} - \Omega P_\varphi = 0, \quad (\text{A3})$$

$$\frac{dP_\varphi}{dt} + \Omega P_r = F, \quad (\text{A4})$$

where the first equation describes dynamics of motion and the second equation defines a value of the reaction force. We have taken into account the relations for unit vectors in polar coordinates: $\frac{d\mathbf{e}_r}{dt} = \Omega \mathbf{e}_\varphi$, $\frac{d\mathbf{e}_\varphi}{dt} = -\Omega \mathbf{e}_r$ and applied the fact that the reaction force is perpendicular to the wire: $\mathbf{F} = \mathbf{e}_\varphi F$.

After quite straightforward calculations, one can show that equation (A3) reduces to (Machabeli & Rogava 1994)

$$\frac{d^2 r}{dt^2} = \frac{\Omega^2 r}{1 - \frac{\Omega^2 r^2}{c^2}} \left[1 - \frac{\Omega^2 r^2}{c^2} - 2 \left(\frac{dr}{dt} \right)^2 \right], \quad (\text{A5})$$

having the following solution

$$r(t) = \frac{v_0}{\Omega} \frac{sn(\Omega t)}{dn(\Omega t)}. \quad (\text{A6})$$

Here, v_0 is the initial radial velocity of the particle and sn and dn are a Jacobian elliptical sine and a modulus, respectively (Abramowitz & Stegun 1965). It is worth noting that for initially relativistic particles ($v_0 \approx c$) the aforementioned expression, approximately reduces to

$$r(t) \approx \frac{v_0}{\Omega} \sin(\Omega t). \quad (\text{A7})$$

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